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DOES TURBULENCE IN THE IRON CONVECTION ZONE CAUSE THE MASSIVE OUTBURSTS OF η CARINAE?

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ABSTRACT

Taken at face value, the observed properties of the central object in η Carinae suggest a very massive, hot main-sequence star, only slightly evolved. If this is so, the star's extraordinarily high steady rate of mass loss must dynamically perturb its outer envelope down to the iron convection zone, in which the kinetic energy associated with turbulent convection can be directly fed into mass ejection. Runaway mass loss, triggered by either internal (pulsational, rotational) or external (tidal) forcing, would produce a secular oscillation of the outer envelope. In either case, the oscillation is potentially able to account for the observed ∼5 yr cycles of visual outbursts in η Car, including the giant eruption of 1843.
Subject headings: convection — stars: individual (η Carinae) — stars: oscillations — stars: variables: other (luminous blue variables) — turbulence

1. INTRODUCTION

One of the brightest nonexplosive stars in the Galaxy is η Carinae. Often classed as a luminous blue variable (LBV), it displays visual outbursts, occasionally of large intensity, at intervals of ∼5 yr (see the recent reviews in Morse, Humphreys, & Damineli 1999; also Davidson & Humphreys 1997; Stothers & Chin 1997). Another oddity of this star is its large surrounding dust nebula, which was ejected in a massive eruption that began in (or shortly before) 1843 and continued for about 20 yr. Despite being still enshrouded by dust, the hot central object is believed to have the following characteristics: luminosity $L = 5 \times 10^6\ L_\odot$ if single or possibly as low as $2.5 \times 10^6\ L_\odot$ if double, effective temperature $T_e = (1.5–4) \times 10^4$ K, and current rate of mass loss $-dM/dt = (0.3–3) \times 10^{-3}\ M_\odot$ yr^{-1}. Although the star's ejecta are clearly enriched in helium and nitrogen, the measured helium abundance, $Y = 0.3–0.4$ (Davidson et al. 1986; Dufour et al. 1999), does not exceed by much the star's natal value, $Y \approx 0.27$. In view of these observations and because the present rate of mass loss (believed to be that of the steady state stellar wind) cannot be sustained for more than $\sim 10^5$ yr without dissipating the star's entire mass, η Car is most likely a very young main-sequence star of 150–300 $M_\odot$ that is evolving quasi-homogeneously because of heavy mass loss (Stothers & Chin 1999). The a priori probability that η Car is a main-sequence object is high because only a few nonexplosive stars with $L > 2 \times 10^6\ L_\odot$ are known and all of them are quite similar in type (Humphreys, Davidson, & Smith 1999).

What, then, could possibly be causing the violent instability? Although many theories have been proposed, none has yet gained general acceptance. Therefore, it seems worth while to adopt, as a working hypothesis, the notion that η Car is, in effect, a single, massive hot main-sequence star early in the phase of core hydrogen burning. Since at least the large eruption of 1843 is considered, also hypothetically, to have originated deep within the outer part of the stellar envelope ($T \sim 10^5$ K or even higher), any hydrodynamical instabilities that are triggered in the hydrogen and helium ionization zones, which for such a hot star must lie near the surface, can probably be counted out as a primary cause.

Damineli (1996) and Damineli, Conti, & Lopes (1997) have suggested that the ∼5 yr cycles in η Car are caused by two colliding stellar winds at periastron in an eccentric double-star system with an orbital period of 5.52 yr. But the massive eruption of 1843 remains unexplained. Since the most recent, and best studied, outburst looks in some ways like an unusual LBV (S Doradus–type) event (Sterken et al. 1999; Davidson et al. 1999; McGregor, Rathborne, & Humphreys 1999), all the outbursts can probably be treated as arising from instabilities deep in the outer envelope of a single star (which may, however, be triggered by periodic tidal forcing by a stellar companion).

Steady mass loss due to the stellar wind in η Car is so large that the star's outer envelope must, from this cause alone, be constantly in a dynamic state (Stothers & Chin 1997). Here, we examine more closely the effect of the enormous stellar wind on the envelope structure. A potential outburst mechanism is found to exist in a large region of turbulent convection that lies deep within the outer envelope.

2. STELLAR MODELS

Ionization of the heaviest elements—especially the most abundant heavy metal, iron—produces a large number of weak spectral absorption lines. These lines dominate the stellar opacity in the temperature range 10^5–10^6 K and furnish two local opacity peaks: a large peak around 2×10^5 K and a smaller one around 1.5×10^6 K (Rogers & Iglesias 1992; Iglesias, Rogers, & Wilson 1992). As the stellar mass increases, envelope densities decline, and, consequently, electron scattering becomes a larger fraction of the total opacity, making the two iron opacity peaks more conspicuous by contrast. Since an increase of stellar mass also raises the contribution of radiation pressure to the total pressure, the envelope becomes less stable against convection. Owing to the main iron opacity peak, convection manages to develop if the stellar mass exceeds $\sim 7\ M_\odot$. Above $\sim 120\ M_\odot$, the smaller opacity peak leads to a second, but much weaker, convection zone. These two convection zones have been termed "iron convection zones" (Stothers & Chin 1993a, 1993b).

Although the derived threshold masses depend somewhat on chemical composition, rotation rate, and evolutionary age, they have been stated here for nonrotating zero-age main-sequence stars with (hydrogen, metals) abundances of $(X, Z) = (0.70,\ 0.03)$. The radius fractions covered by the two iron convection zones and by the convective core are shown as a func-

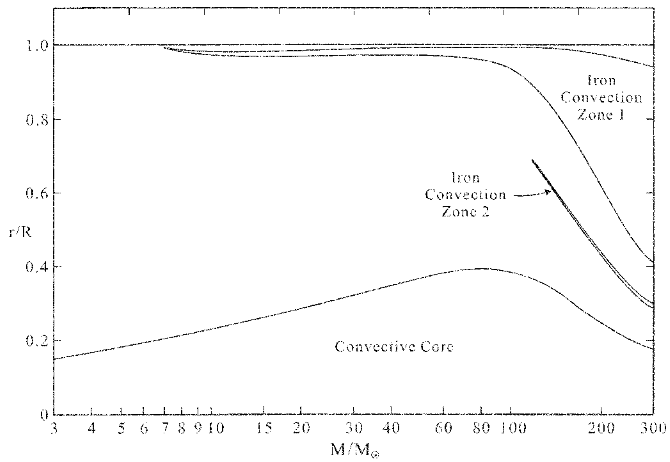

FIG. 1.—Spatial extent of the convective zones in homogenous hydrogen-burning stellar models with $(X, Z) = (0.70, 0.03)$. Note that the radius fraction contained in the convective core is not monotonic with increasing stellar mass (however, the mass fraction is). The mass fractions of the outer convective zones are shown in Figs. 2 and 3.

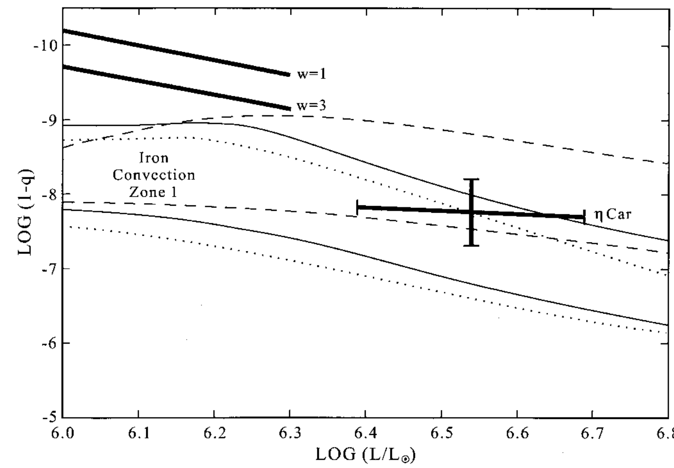

FIG. 2.—Logarithm of the mass depth below the surface as a function of the logarithm of the stellar luminosity. Locations of the upper and lower boundaries of the main iron convection zone are shown for $(X, Z) =$ (0.70, 0.03) (*solid lines*), (0.70, 0.02) (*dashed lines*), and (0.35, 0.03) (*dotted lines*). The bottom of the dynamically perturbed region is indicated in the case of ordinary O stars with $L < 2 \times 10^6\ L_\odot$ ($w = 1$ and $w = 3$) and η Car, assumed in all cases to be nonrotating.

tion of stellar mass in Figure 1. The models strongly resemble those calculated with the large (although erroneous) Carson (1976) opacities, which yielded a large convection zone in the envelope at roughly the same temperatures (Fig. 3 of Stothers 1976). However, our new models are not very sensitive to the choice of convective mixing length, which is here taken to be 1.4 times the local pressure scale height.

To study the effects of a lower metallicity and of an evolutionary helium enrichment, we set also $(X, Z) =$ (0.70, 0.02) and (0.35, 0.03). The second choice simulates an intermediate stage in the quasi-homogeneous evolution of a very massive star, which can be represented fairly well by stellar models with a homogeneous composition for values of X that are not too small (Stothers & Chin 1999). Uniform rotation is treated by employing the rotational equations used by Sackmann & Anand (1970) and applying them to the limiting case of rotational breakup at the equator.

In the case of main-sequence stars with luminosities up to $2 \times 10^6\ L_\odot$, mass loss probably has no important dynamical effect on the outer envelope. To justify this assertion, we use in the stellar models the standard mass-loss rates of Nieuwenhuijzen & de Jager (1990), which have recently been supported by the work of Lamers & Leitherer (1993) and Puls et al. (1996). Because, however, de Koter, Heap, & Hubeny (1997) and Crowther & Bohannan (1997) have derived rates that are ~3 times larger for the most luminous O stars, we shall consider also the case in which the standard rates are multiplied by a factor $w = 3$. For η Car we adopt $-dM/dt = 1 \times 10^{-3}\ M_\odot\ \mathrm{yr}^{-1}$, which has an estimated possible error of a factor 3 (Davidson et al. 1986, 1995; White et al. 1994; Cox et al. 1995; Hillier 1999).

3. DISCUSSION

The models calculated here are fully hydrostatic, and therefore any quantities derived from them can only approximately represent the structure of real stars that suffer mass loss heavy enough to affect the outer layers. Nevertheless, we suppose that these models provide a sufficiently accurate zeroth-order approximation for our present purposes. The depth to which surface mass loss directly affects the envelope structure can be represented as a perturbed fraction of the total mass,

$$\delta M/M = |dM/dt|\tau_{\rm dyn}/M, \qquad (1)$$

where $\tau_{\rm dyn}$ is the dynamical timescale of the outer envelope. Since $\tau_{\rm dyn}$ equals simply the period Π of the fundamental mode of radial pulsation, for which the nondimensional eigenvalue $\omega^2 = (2\pi/\Pi)^2 R^3/GM$ lies very close to 3 in the case of very massive main-sequence stars (Stothers 1992), it becomes a simple matter to compute $\delta M/M$ for any assigned rate of mass loss, $-dM/dt$, by substituting Π for $\tau_{\rm dyn}$. Then the bottom of the perturbed region occurs at a mass fraction $(1-q)_{\rm dyn} = \delta M/M$ below the surface, where $q = M(r)/M$. For η Car, Π is ~1 day.

Figures 2 and 3 show, as a function of stellar luminosity,

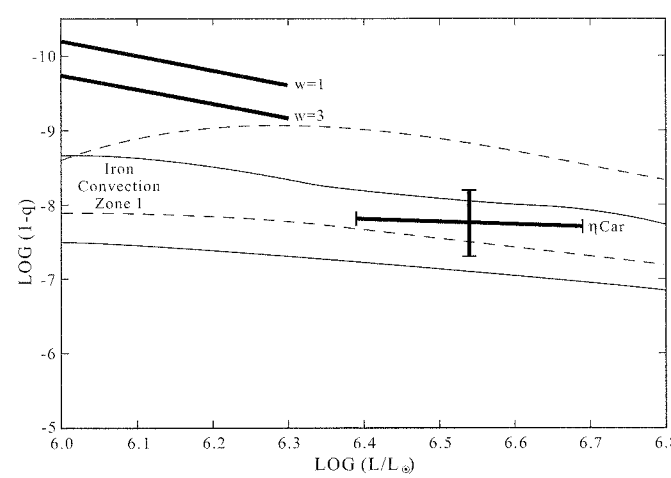

FIG. 3.—Logarithm of the mass depth below the surface as a function of the logarithm of the stellar luminosity. Locations of the upper and lower boundaries of the main iron convection zone are shown for nonrotating stars (*dashed lines*) and uniformly rotating stars at breakup (*solid lines*), with $(X, Z) = (0.70, 0.02)$. The bottom of the dynamically perturbed region is indicated in the case of ordinary O stars with $L < 2 \times 10^6\ L_\odot$ ($w = 1$ and $w = 3$) and η Car.

the perturbed depth $(1 - q)_{\rm dyn}$ for $w = 1$, $w = 3$, and $-dM/dt = 1 \times 10^{-3}\ M_\odot\ {\rm yr}^{-1}$. Also plotted are the upper and lower boundaries of the main iron convection zone. Notice that Figure 2 includes the whole range of likely chemical compositions for η Car, while Figure 3 covers all possible rates of uniform rotation. Although helium enrichment proves to make little difference, thus removing evolution as an important factor for consideration, there do exist modest increases of the size of the iron convection zone for greater metallicities and for faster rotation rates.

The surface wind from an ordinary O star with $L < 2 \times 10^6\ L_\odot$ perturbs only a tiny fraction of the star's mass, all of it in the radiative, outer part of the envelope. On the other hand, the stellar wind's influence in η Car extends down to $1 - q = 10^{-8}$, which falls in the middle of the iron convection zone. This difference is crucial. All other types of stars that possess convective envelopes generate their winds in or near their photospheres, where the convective flux is very weak. In η Car, however, the massive wind dissipates the outer layers so fast that the bottom of the wind taps directly into a highly energetic field of large-scale turbulence.

Assuming the adequacy of classical mixing-length theory for estimating the gross properties of the turbulent convective flows (Tennekes & Lumley 1972), our present models predict local mean convective velocities attaining ~0.4 times sound speed throughout the iron convection zone. This means that the convective overturning time must be of the order of $\tau_{\rm dyn}$. Since the convective heat flux reaches 0.3 of the total flux and the turbulent energy flux accounts for an additional 0.01 of the total, a large reservoir of kinetic energy is directly available to be fed into mass loss. It seems a plausible assumption, akin to one made by Forbes (1968) for red giants, to set the mass-loss rate at the stellar surface equal to a numerical constant times the mass of that part of the iron convection zone that lies within the perturbed region of the star. Roughly then,

$$-dM/dt = K(1 - q)_{\rm dyn}, \qquad (2)$$

where K is a constant. This assumption has some justification in that equation (2) is then compatible with equation (1).

Let us now suppose that as a result of some fluctuation of the turbulent flow the mass-loss rate begins to grow. The increased mass-loss rate will perturb more of the iron convection zone, thereby further enhancing the mass outflow according to equation (2). As matter streams from the surface of the star, new material from below rises to replace it. If mass loss continues long enough, the equivalent of many equilibrium outer envelopes will be removed. This runaway process ends either when the base of the perturbed region penetrates down into the underlying stable radiative layers or when the initial triggering disturbance dies out. Thereafter, the mass-loss rate should return to something like its original value.

From Figures 2 and 3 we readily predict that, to have a modest outburst, only a small increase of the present rate of mass loss from η Car is needed. On the other hand, the maximum possible rate consistent with our stellar models corresponds to an increase of $(1 - q)_{\rm dyn}$ by a factor of between 10 and 10^2, putting the base of the outer envelope well into the radiative layers; the rate must then have reached 10^{-2} to $10^{-1}\ M_\odot\ {\rm yr}^{-1}$. This lies in the range of η Car's probable rate during the ~20 yr period following 1843, which is variously estimated to have been $4 \times 10^{-3}\ M_\odot\ {\rm yr}^{-1}$ (van Genderen & Thé 1984), $2 \times 10^{-2}\ M_\odot\ {\rm yr}^{-1}$ (Hyland et al. 1979), $7.5 \times 10^{-2}\ M_\odot\ {\rm yr}^{-1}$ (Andriesse, Donn, & Viotti 1978), and $\sim 10^{-1}\ M_\odot\ {\rm yr}^{-1}$ (Davidson 1989).

The runaway process is expected to repeat cyclically. The mean cycle time for arbitrarily large mass-loss rates has already been determined for massive main-sequence stars from quasi-static evolutionary models and is ~4 yr for the luminosity of η Car (Stothers & Chin 1997). This value essentially equals the thermal relaxation time of the outer envelope ($\tau_{\rm th} = E_{\rm th}/L$) and agrees well with the ~5 yr cycle time shown by η Car. Since it is also close to the 5.52 yr orbital period proposed by Damineli (1996), tidal forcing by a stellar companion (if such a companion exists) might be both tuning and regularizing what might otherwise be a quasi-regular cyclicity of the type displayed by the classical LBVs.

If our model is correct, the great eruption in 1843 must have been triggered by considerably larger than average turbulent fluctuations in the iron convection zone. What actually amplified the fluctuations is conjectural, although pulsational-mode interactions, angular momentum adjustments, or even tidal forcing may have been at work. At the peak of the eruption, our zeroth-order approximation for the structure of the outer envelope must break down near the surface. The huge mass-loss rate requires an energy supply rate

$$\delta L = \left(\frac{GM}{R} + \frac{1}{2}v_\infty^2\right)\left|\frac{dM}{dt}\right|, \qquad (3)$$

where v_∞ is the terminal velocity of the ejected matter (Forbes 1968). For η Car we adopt $v_\infty = 1 \times 10^3\ {\rm km\ s^{-1}}$ (Walborn, Blanco, & Thackeray 1978; Meaburn et al. 1993), $M = 250\ M_\odot$, $R = 60\ R_\odot$, and $L = 5 \times 10^6\ L_\odot$ (Stothers & Chin 1997, 1999). If the mass-loss rate actually attained $\sim 10^{-1}\ M_\odot\ {\rm yr}^{-1}$ during the great eruption (Davidson 1989), then δL must have been $\sim L$. This large energy drain would have caused an equivalent drop in luminosity at the stellar surface. At the base of the outer envelope where the mass loss begins, the luminosity drop would have been close to zero because the interior layers of the star would still have been transporting the star's equilibrium luminosity.

In 1843, however, a drastic surface luminosity drop was not observed. In fact, there was probably a bolometric brightening by ~2.5 mag (van Genderen & Thé 1984). Davidson (1989) has pointed out that the ejected material, once it was liberated from the star, would have radiated away its substantial internal energy. Since the total energy of the system (star plus ejecta) must be conserved over a short time interval, this extra radiated energy would have equaled, almost exactly, the energy losses sustained by the star's interior.

Hydrodynamical modeling of so complex and extreme an object as η Car, involving detailed time-dependent interactions between mass loss and turbulence (to say nothing of pulsation, rotation, and tides), is not yet feasible, but the arguments made here suggest a model of η Car that may explain its instability in what is possibly the most straightforward way.

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REFERENCES

Andriesse, C. D., Donn, B. D., & Viotti, R. 1978, MNRAS, 185, 771
Carson, T. R. 1976, ARA&A, 14, 95
Cox, P., et al. 1995, A&A, 297, 168
Crowther, P. A., & Bohannan, B. 1997, A&A, 317, 532
Damineli, A. 1996, ApJ, 460, L49
Damineli, A., Conti, P., & Lopes, D. F. 1997, NewA, 2, 107
Davidson, K. 1989, in Physics of Luminous Blue Variables, ed. K. Davidson, A. F. J. Moffat, & H. J. G. L. M. Lamers (Dordrecht: Kluwer), 101
Davidson, K., Dufour, R. J., Walborn, N. R., & Gull, T. R. 1986, ApJ, 305, 867
Davidson, K., & Humphreys, R. M. 1997, ARA&A, 35, 1
Davidson, K., Ishibashi, K., Gull, T. R., & Humphreys, R. M. 1999, in ASP Conf. Ser. 179, Eta Carinae at the Millenium, ed. J. A. Morse, R. M. Humphreys, & A. Damineli (San Francisco: ASP), 227
Davidson, K., et al. 1995, AJ, 109, 1784
de Koter, A., Heap, S. R., & Hubeny, I. 1997, ApJ, 477, 792
Dufour, R. J., Glover, T. W., Hester, J. J., Currie, D. G., van Orsow, D., & Walter, D. K. 1999, in ASP Conf. Ser. 179, Eta Carinae at the Millenium, ed. J. A. Morse, R. M. Humphreys, & A. Damineli (San Francisco: ASP), 134
Forbes, J. E. 1968, ApJ, 153, 495
Hillier, D. J. 1999, in ASP Conf. Ser. 179, Eta Carinae at the Millenium, ed. J. A. Morse, R. M. Humphreys, & A. Damineli (San Francisco: ASP), 98
Humphreys, R. M., Davidson, K., & Smith, N. 1999, PASP, 111, 1124
Hyland, A. R., Robinson, G., Mitchell, R. M., Thomas, J. A., & Becklin, E. E. 1979, ApJ, 233, 145
Iglesias, C. A., Rogers, F. J., & Wilson, B. G. 1992, ApJ, 397, 717
Lamers, H. J. G. L. M., & Leitherer, C. 1993, ApJ, 412, 771
McGregor, P. J., Rathborne, J. M., & Humphreys, R. M. 1999, in ASP Conf. Ser. 179, Eta Carinae at the Millenium, ed. J. A. Morse, R. M. Humphreys, & A. Damineli (San Francisco: ASP), 236
Meaburn, J., Gehring, G., Walsh, J. R., Palmer, J. W., Lopez, J. A., Bryce, M., & Raga, A. C. 1993, A&A, 276, L21
Morse, J. A., Humphreys, R. M., & Damineli, A., eds. 1999, ASP Conf. Ser. 179, Eta Carinae at the Millenium (San Francisco: ASP)
Nieuwenhuijzen, H., & de Jager, C. 1990, A&A, 231, 134
Puls, J., et al. 1996, A&A, 305, 171
Rogers, F. J., & Iglesias, C. A. 1992, ApJS, 79, 507
Sackmann, I. J., & Anand, S. P. S. 1970, ApJ, 162, 105
Sterken, C., Freyhammer, L., Arentoft, T., & van Genderen, A. M. 1999, A&A, 346, L33
Stothers, R. B. 1976, ApJ, 209, 800
———. 1992, ApJ, 392, 706
Stothers, R. B., & Chin, C.-w. 1993a, ApJ, 408, L85
———. 1993b, ApJ, 412, 294
———. 1997, ApJ, 489, 319
———. 1999, ApJ, 522, 960
Tennekes, H., & Lumley, J. L. 1972, A First Course in Turbulence (Cambridge: MIT Press)
van Genderen, A. M., & Thé, P. S. 1984, Space Sci. Rev., 39, 317
Walborn, N. R., Blanco, B. M., & Thackeray, A. D. 1978, ApJ, 219, 498
White, S. M., Duncan, R. A., Lim, J., Nelson, G. J., Drake, S. A., & Kundu, M. R. 1994, ApJ, 429, 380